The scale of supersymmetry breaking as a free parameter

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While supersymmetric extensions of the Standard Model can be fully described in terms of explicitly broken global supersymmetry, this description is only effective. Once related to spontaneous breaking in a more fundamental theory, the effective parameters translate to functions of two distinct scales, the scale of spontaneous supersymmetry breaking and the scale of its mediation to the standard-model fields. The scale dependence will be written explicitly and the full spectrum of supersymmetry breaking operators which emerges will be explored. It will be shown that, contrary to common lore, scale-dependent operators can play an important role in determining the phenomenology. For example, theories with low-energy supersymmetry breaking, such as gauge mediation, may correspond to a scalar potential which is quite different than in theories with high-energy supersymmetry breaking, such as gravity mediation. As a concrete example, the Higgs mass prediction will be discussed in some detail and its upper bound will be shown to be sensitive to the supersymmetry breaking scale.

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1. INTRODUCTION

If supersymmetry is realized in nature, and if it plays a role in resolving the hierarchy problem associated with the quadratically divergent quantum corrections to the Standard Model (SM) Higgs boson, then it is explicitly broken (as dictated by experiment) at the electroweak scale (as dictated by its solution to the hierarchy problem). The breaking is often assumed to be soft so that the theory is at most logarithmically divergent and no new hierarchy problem appears. Hence, the dynamics of the scalar fields ϕ, for example, is described by a potential of the form

$$\begin{aligned} V &= \left|\frac{\partial W}{\partial \Phi}\right|^2 + \frac{g^2}{2}\left|\sum_i \phi_i T_i \phi_i^\dagger\right|^2 + m^2\phi\phi^\dagger \\ &+ (B\phi^2 + A\phi^3 + A'\phi^2\phi^\dagger + \text{H.c.}). \qquad (1) \end{aligned}$$

In the first (F-)term, W is the superpotential,[2] $W(\Phi) = \mu H_1 H_2 +$ Yukawa terms, describing the interactions of the matter and Higgs fields, with H_1 (H_2) being the negative- (positive-)hypercharge Higgs doublet. In the second (D-)term, g is the gauge coupling, and a summation over all gauge groups and generators T is implied. The $F-$ and $D-$terms dictate, in this description, all quartic interactions which preserve (up to logarithmic corrections [1]) the supersymmetric structure. All other terms break supersymmetry softly. Note that the soft supersymmetry breaking (SSB) parameters are necessarily dimensionful (due to the softness requirement).

Indeed, this is the minimal description of supersymmetric extensions of the SM, with the exception of the non-holomorphic A'-parameters which may not be soft in models containing pure singlets (and which are often omitted). While the $F-$ and $D-$terms preserve supersymmetry and depend on the gauge and Yukawa structure of the low-energy theory, the SSB terms should depend on the ultraviolet supersymmetry breaking vacuum expectation values (VEVs). As such, they depend on, and encode, hidden ultraviolet physics.

Spontaneous supersymmetry breaking is most conveniently parameterized in terms of a spurion

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[2] $\Phi = \phi + \theta\psi$ is a chiral superfield, θ is the superspace coordinate, and ψ is a fermion.

field $X = \theta^2 F$ with a non vanishing F-VEV, which is an order parameter of (global) supersymmetry breaking. All SSB parameters can be written as nonrenormalizable operators which couple the spurion X to the SM superfields. The operators are suppressed by the scale of the mediation of supersymmetry breaking from the (hidden) sector parameterized by the spurion to the SM (observable) sector. The coefficients of the various operators are dictated by the nature of the interaction between the sectors and by the loop-order at which it occurs. In the following, we will explicitly write all operators which couple the two sectors. This will allow us to consider a more general and complete form of the supersymmetry breaking potential, which still resolves the hierarchy problem. While we will also identify the operators which correspond to the SSB parameters, our focus will be on those operators which are not included in the minimal form of the potential (1). This will be done in Sec. 2.

Eventhough the SSB parameters appearing in (1) are functions of the relevant ultraviolet scales, their magnitude is uniquely determined by the assumption that supersymmetry stabilizes the weak scale against divergent quantum corrections, $m \sim \mathcal{O}(M_{\rm Weak}) \sim \mathcal{O}(100)$ GeV etc. This places a constraint on the ratio of the supersymmetry breaking scale $\sqrt{F}$ and the scale of its mediation M such that $F/M \simeq M_{\rm Weak}$, and it implies that no information can be extracted[3] regarding either scale from the SSB parameters. (We will comment on the case of A' below.) The generalization of (1) introduces dimensionless hard supersymmetry breaking (HSB) parameters in the potential. Their magnitude depends on the various scales in a way which both provides useful information (unlike the SSB parameters) and does not destabilize the solution to the hierarchy problem. In Sec. 3 we will demonstrate this while calculating the Higgs mass, which will provide a concrete example. We conclude in Sec. 4.

[3]Specific ultraviolet relations between the various SSB parameters could still be studied using the renormalization group formalism.

2. CLASSIFICATION OF OPERATORS

It is convenient to contain, without loss of generality, all observable-hidden interactions in the non-holomorphic Kahler potential K, $\mathcal{L} = \int d^2\theta d^2\bar{\theta}$ (which is not protected by nonrenormalization theorems). We therefore turn to a general classification of K-operators, originally presented in Ref. [2]. (See also Ref. [3].) We do not impose any global symmetries, which can obviously eliminate some of the operators, and we keep all operators which survive the superspace integration. (Note that the spurion is assumed to have only a F-VEV.)

The effective low-energy Kahler potential of a rigid $N = 1$ supersymmetry theory[4] is given by

$$\begin{aligned} K &= K_0(X, X^\dagger) + K_0(\Phi, \Phi^\dagger) \\ &+ \frac{1}{M} K_1(X, X^\dagger, \Phi, \Phi^\dagger) \\ &+ \frac{1}{M^2} K_2(X, X^\dagger, \Phi, \Phi^\dagger) \\ &+ \frac{1}{M^3} K_3(X, X^\dagger, \Phi, \Phi^\dagger, D_\alpha, W_\alpha) \\ &+ \frac{1}{M^4} K_4(X, X^\dagger, \Phi, \Phi^\dagger, D_\alpha, W_\alpha) + \cdots \end{aligned} \tag{2}$$

where, as before, X is the spurion and Φ are the chiral superfields of the low-energy theory. D_α is the covariant derivative with respect to the superspace chiral coordinate θ_α, and W_α is the $N = 1$ gauge supermultiplet in its chiral representation, $W_\alpha \sim \lambda_\alpha + \theta_\alpha V$, with λ and V here being the gaugino and vector boson, respectively. Once a separation between supersymmetry breaking field X and low-energy Φ fields is imposed, there is no tree-level renormalizable interaction between the two sets of fields, and their mixing can arise only at the non-renormalizable level $K_{l\geq 1}$.

The superspace integration $\mathcal{L}_D = \int d^2\theta d^2\bar{\theta} K$ reduces K_1 and K_2 to the usual SSB terms, as well as the superpotential μ-parameter $W \sim \mu\Phi^2$, which were discussed in the previous section. It also contains Yukawa operators $W \sim y\Phi^3$ which can appear in the effective low-energy superpotential. These are summarized in Tables 1 and 2. (We did not include linear terms that may appear

[4]This description is useful also to $N = 2$ supersymmetry under certain assumptions [2].

in the case of a singlet superfield.) Finally, The last term in Table 1 contains correlated but unusual quartic and Yukawa couplings. They are soft as they involve at most logarithmic divergences.

Integration over K_3 produces non-standard soft terms, one of which was discussed in the previous section. These terms are soft unless the theory contains a pure singlet field, in which case they can induce a quadratically divergent linear term. They are summarized in Table 3. The integration over K_3 also generates contributions to the ("standard") $A-$ and gaugino-mass terms. These terms could arise at lower orders in $\sqrt{F}/M$ from integration over holomorphic functions (and in the case of A, also from K_1). However, this is equivalent to integration over K if $\int d^2\bar{\theta}(X^\dagger/M^2) \simeq 1$. Note that in the presence of superpotential Yukawa couplings, a supersymmetry breaking Higgsino mass term $\tilde{\mu}\widetilde{H}_1\widetilde{H}_2$ can be rotated to a combination of μ and A' terms, and vice versa.

Lastly, superspace integration over K_4 leads to dimensionless hard operators. These are summarized in Table 4, and will occupy the remaining of this lecture. Table 4 also contain supersymmetry breaking gauge-Yukawa interactions $\sim \phi^*\psi\lambda$. This is equivalent to the HSB kinetic term for the gauginos which was discussed recently in Ref. [4]. (HSB gaugino couplings are also generated radiatively in the presence of SSB [1,5].)

Higher orders in $(1/M)$ can be safely neglected as supersymmetry and the superspace integration allow only a finite expansion in $\sqrt{F_X}/M$, that is $\mathcal{L} = f[F_X^n/M^l]$ with $n \leq 2$ and l is the index K_l in expansion Eq. (2). Hence, terms with $l > 4$ are suppressed by at least $(\langle X\rangle/M)^{l-4}$. We will assume the limit $\langle X\rangle \ll M$ for the supersymmetry preserving VEV $\langle X\rangle$, i.e., $X \sim \theta^2 F_X$, so that all such operators can indeed be neglected and the expansion is rendered finite.

It is interesting to identify two phenomenologically interesting groups of terms in K, (i) those terms which can break the chiral symmetries and can generate Yukawa terms in the low-energy effective theory, and (ii) new sources for quartic interactions.

The relevant chiral symmetry breaking terms

Table 1
The soft supersymmetry breaking terms as operators contained in K_1 and K_2. $\Phi = \phi + \theta\psi + \theta^2 F$ is a low-energy superfield while X, $\langle F_X\rangle \neq 0$, parameterizes supersymmetry breaking. $F^\dagger = \partial W/\partial\Phi$. The infrared operators are obtained by superspace integration over the ultraviolet operators.

ultraviolet K operator	infrared $\mathcal{L}_D$ operator
$\frac{X}{M}\Phi\Phi^\dagger + \text{H.c.}$	$A\phi F^\dagger + \text{H.c.}$
$\frac{XX^\dagger}{M^2}\Phi\Phi^\dagger + \text{H.c.}$	$\frac{m^2}{2}\phi\phi^\dagger + \text{H.c.}$
$\frac{XX^\dagger}{M^2}\Phi\Phi + \text{H.c.}$	$B\phi\phi + \text{H.c.}$
$\frac{X^\dagger}{M^2}\Phi^2\Phi^\dagger + \text{H.c.}$	$\kappa\phi^\dagger\phi F + \text{H.c.}$ $y\phi^\dagger\psi\psi + \text{H.c.}$

Table 2
The effective renormalizable $N = 1$ superpotential W operators contained in K_1 and K_2, $\mathcal{L} = \int d^2\theta W$. Symbols are defined in Table 1. The infrared operators are obtained by superspace integration over the ultraviolet operators.

ultraviolet K operator	infrared W operator
$\frac{X^\dagger}{M}\Phi^2 + \text{H.c.}$	$\mu\Phi^2$
$\frac{X^\dagger}{M^2}\Phi^3 + \text{H.c.}$	$y\Phi^3$

Table 3
The non-standard or semi-hard supersymmetry breaking terms as operators contained in K_3. W^α is the $N=1$ chiral representation of the gauge supermultiplet and λ is the respective gaugino. D_α is the covariant derivative with respect to the (explicit) superspace coordinate θ_α. All other symbols are as in Table 1. The infrared operators are obtained by superspace integration over the ultraviolet operators.

ultraviolet K operator	infrared $\mathcal{L}_D$ operator
$\frac{XX^\dagger}{M^3}\Phi^3 + \text{H.c.}$	$A\phi^3 + \text{H.c.}$
$\frac{XX^\dagger}{M^3}\Phi^2\Phi^\dagger + \text{H.c.}$	$A'\phi^2\phi^\dagger + \text{H.c.}$
$\frac{XX^\dagger}{M^3}D^\alpha\Phi D_\alpha\Phi + \text{H.c.}$	$\tilde{\mu}\psi\psi + \text{H.c.}$
$\frac{XX^\dagger}{M^3}D^\alpha\Phi W_\alpha + \text{H.c.}$	$M'_\lambda\psi\lambda + \text{H.c.}$
$\frac{XX^\dagger}{M^3}W^\alpha W_\alpha + \text{H.c.}$	$\frac{M_\lambda}{2}\lambda\lambda + \text{H.c.}$

in tables 1 and 3 can be identified with $A-$ and $A'-$terms which couple the matter sfermions to the Higgs fields of electroweak symmetry breaking. The chiral symmetry breaking originates in this case in the scalar potential and propagates to the fermions at one loop [6]. More interestingly, a generic Kahler potential is also found to contain tree-level chiral Yukawa couplings. These include $\mathcal{O}(F_X/M^2)$ supersymmetry conserving and soft couplings and $\mathcal{O}(F_X^2/M^4)$ hard chiral symmetry breaking couplings, leading to new avenues for fermion mass generation [2].

Quartic coupling arise at $\mathcal{O}(F_X/M^2)$, from supersymmetry conserving operators in Table 1 (depending on F_Φ), and at $\mathcal{O}(F_X^2/M^4)$ from hard couplings in Table 4. They can potentially alter the supersymmetry conserving nature of the quartic potential in (1).

The relative importance of the HSB operators relates to a more fundamental question: What are the scales $\sqrt{F_X}$ and M? This will be addressed in Sec. 3. However, before doing so we need to address a different question regarding the

Table 4
The dimensionless hard supersymmetry breaking terms as operators contained in K_4. Symbols are defined as in Tables 1 and 3. The infrared operators are obtained by superspace integration over the ultraviolet operators.

ultraviolet K operator	infrared $\mathcal{L}_D$ operator
$\frac{XX^\dagger}{M^4}\Phi D^\alpha\Phi D_\alpha\Phi + \text{H.c.}$	$y\phi\psi\psi + \text{H.c.}$
$\frac{XX^\dagger}{M^4}\Phi^\dagger D^\alpha\Phi D_\alpha\Phi + \text{H.c.}$	$y\phi^\dagger\psi\psi + \text{H.c.}$
$\frac{XX^\dagger}{M^4}\Phi D^\alpha\Phi W_\alpha + \text{H.c.}$	$\bar{y}\phi\psi\lambda + \text{H.c.}$
$\frac{XX^\dagger}{M^4}\Phi^\dagger D^\alpha\Phi W_\alpha + \text{H.c.}$	$\bar{y}\phi^\dagger\psi\lambda + \text{H.c.}$
$\frac{XX^\dagger}{M^4}\Phi W^\alpha W_\alpha + \text{H.c.}$	$\bar{y}\phi\lambda\lambda + \text{H.c.}$
$\frac{XX^\dagger}{M^4}\Phi^\dagger W^\alpha W_\alpha + \text{H.c.}$	$\bar{y}\phi^\dagger\lambda\lambda + \text{H.c.}$
$\frac{XX^\dagger}{M^4}\Phi^2\Phi^{\dagger\,2} + \text{H.c.}$	$\kappa(\phi\phi^\dagger)^2 + \text{H.c.}$
$\frac{XX^\dagger}{M^4}\Phi^3\Phi^\dagger + \text{H.c.}$	$\kappa\phi^3\phi^\dagger + \text{H.c.}$

potentially destabilizing properties of the different HSB operators, which relates to the nature of the cut-off scale Λ. Indeed, one has to confirm that a given theory is not destabilized when the hard operators are included, an issue which is interestingly model independent. In order to do so, consider the implications of the hardness of the operators contained in K_4. Yukawa and quartic couplings can destabilize the scalar potential by corrections Δm^2 to the mass terms of the order of

$$\Delta m^2 \sim \begin{cases} \frac{\kappa}{16\pi^2}\Lambda^2 \sim \frac{1}{16\pi^2}\frac{F_X^2}{M^4}\Lambda^2 \sim \frac{1}{16\pi^2 c_m}m^2 \\ \frac{y^2}{16\pi^2}\Lambda^2 \sim \frac{1}{16\pi^2}\frac{F_X^4}{M^8}\Lambda^2 \sim \frac{1}{16\pi^2 c_m}m^2\frac{m^2}{M^2}, \end{cases} \tag{3}$$

where we identified $\Lambda \simeq M$ and c_m is a dimensionless coefficient omitted in Table 1, $m^2/2 = c_m F_X^2/M^2$. The hard operators were substituted by the appropriate powers of F_X/M^2. Once M is identified as the cut-off scale above which the full supersymmetry is restored, then these terms are harmless as the contributions are bound from above by the tree-level scalar squared-mass parameters. In particular, the softness assumption imposed on the supersymmetry breaking terms in (1) was not necessary. (This observation extends to the case of non-standard soft operators such as $A' \sim F_X^2/M^3$ in the presence of a singlet).

In fact, such hard divergent corrections are well known in $N = 1$ supergravity with $\Lambda = M = M_{\text{Planck}}$, where they perturb any given set of tree-level boundary conditions for the SSB parameters [7]. Given the supersymmetry breaking scale in this case, $F \simeq M_{\text{Weak}}M_{\text{Planck}}$, the Yukawa (and quartic) operators listed below are proportional in these theories to $(M_{\text{Weak}}/M_{\text{Planck}})^n$, $n = 1, 2$, and are often omitted. Nevertheless, such terms can shift any boundary conditions for the SSB by $\mathcal{O}(100\%)$ [7] due to quadratically divergent corrections.

We conclude that, in general, quartic couplings and chiral Yukawa couplings appear once supersymmetry is broken, and if supersymmetry is broken at low energy then these couplings could be sizable yet harmless. We will explore possible implications of the HSB quartic couplings in the next section.

3. THE HIGGS MASS vs. THE SCALE OF SUPERSYMMETRY BREAKING

In the previous section we have shown that, in general, HSB quartic couplings κ_{hard} arise in the scalar potential (from non-renormalizable operators in the Kahler potential, for example). Assuming that the SSB parameters are characterized by a parameter $m_0 \sim 1$ TeV then

$$\kappa_{hard} = \tilde{\kappa}_h \frac{F^2}{M^4} \simeq \tilde{\kappa}_h (16\pi^2)^{2n} \left(\frac{m_0}{M}\right)^2, \tag{4}$$

where M is a dynamically determined scale parameterizing the communication of supersymmetry breaking to the SM sector, which is distinct from the supersymmetry breaking scale $\sqrt{F} \simeq (4\pi)^n\sqrt{m_0 M}$. The exponent $2n$ is the loop order at which the mediation of supersymmetry breaking to the (quadratic) scalar potential occurs. (Non-perturbative dynmaics may lead to different relations that can be described instead by an effective value of n.) The coupling $\tilde{\kappa}_h$ is an unknown dimensionless coupling (for example, in the Kahler potential). As long as such quartic couplings are not arbitrary but are related to the source of the SSB parameters and are therefore described by (4), then they do not destabilize the scalar potential and do not introduce quadratic dependence on the ultraviolet cut-off scale, which is identified with M. Stability of the scalar potential only constrains $\tilde{\kappa}_h \lesssim \min\left((1/16\pi^2)^{2n-1}, 1\right)$ (though calculability and predictability are diminished).

The $F-$ and $D-$term-induced quartic potential in (1) gives for the (pure D-induced) tree-level Higgs coupling, $V = \kappa h^4$,

$$\kappa = \frac{g'^2 + g^2}{4}\cos^2 2\beta, \tag{5}$$

where $\tan\beta \equiv \langle H_2\rangle/\langle H_1\rangle$, and we work in the decoupling in which one physical Higgs doublet H is sufficiently heavy and decouples from electroweak symmetry breaking while a second SM-like Higgs doublet is roughly given by $h \simeq H_1\cos\beta + H_2\sin\beta$, and we conveniently defined

$$H_n = \begin{pmatrix} H_n^+ \\ (H_n^0 + iA_n^0)/\sqrt{2} \end{pmatrix}. \tag{6}$$

The HSB coupling corrects this relation. Given the strict tree-level upper bound that follows from (5), $m_{h^0}^2 \leq M_Z^2 \cos^2 2\beta$, it is suggestive that HSB may not be only encoded in, but also measured via, the Higgs mass. We will explore this possibility, originally pointed out and studied in Ref. [8], in this section.

In the case that supergravity interactions mediate supersymmetry breaking from some "hidden" sector (where supersymmetry is broken spontaneously) to the SM sector, one has $M = M_{\rm Planck}$. The corrections are therefore negligible whether the mediation occurs at tree level ($n = 0$) or loop level ($n \geq 1$) and can be ignored for most purposes. (For exceptions, see Refs. [3,7].) In general, however, the scale of supersymmetry breaking is an arbitrary parameter and depends on the dynamics that mediate the SSB parameters. For example, it was shown recently that in the case of $N = 2$ supersymmetry one expects $M \sim 1\,$TeV [2]. Also, in models with extra large dimensions the fundamental $M_{\rm Planck}$ scale can be as low as a few TeV, leading again to $M \sim 1\,$TeV. (For example, see Ref. [9].) A "TeV-type" mediation scale implies a similar supesymmetry breaking scale and provides an unconventional possibility. (For a discussion, see Ref. [2].) If indeed $M \sim 1\,$TeV then κ_{hard} given in (4) is $\mathcal{O}(1)$ (assuming tree-level mediation (TLM) and $\mathcal{O}(1)$ couplings $\tilde{\kappa}_h$ in the Kahler potential). The effects on the Higgs mass must be considered in this case.

Though one may argue that TLM models represent a theoretical extreme, this is definitely a viable possibility. A more familiar and surprising example is given by the (low-energy) gauge mediation (GM) framework [10]. In GM, SM gauge loops communicate between the SM fields and some messenger sectors, mediating the SSB potential. The Higgs sector and the related operators, however, are poorly understood in this framework [11] and therefore all allowed operators should be considered. In its minimal incarnation (MGM) $2n = 2$, and $M \sim 16\pi^2 m_0 \sim 100\,$TeV parameterizes both the mediation and supersymmetry breaking scales. The constraint (3) corresponds to $\kappa_{hard} \sim \tilde{\kappa}_h \lesssim 1/16\pi^2$ and the contribution of $\delta\kappa_{hard}$ to the Higgs mass could be comparable to the contribution of the supersymmetric coupling (5). A particularly interesting case is that of non-perturbative messenger dynamics (NPGM) in which case $n_{eff} = 1/2$, $M \sim 4\pi m_0 \sim 10\,$TeV [12], and the constraint on $\tilde{\kappa}_h$ is relaxed to 1. Now $\kappa_{hard} \lesssim 1$ terms could dominate the Higgs mass. The various frameworks are summarized in Table 5.

In order to address the β-dependence of the HSB contributions (which is different from that of all other terms) we recall the general two-Higgs-doublet model (2HDM)[13]. The Higgs quartic potential can be written down as[5]

$$\begin{aligned} V_{\phi^4} &= \frac{1}{2}\kappa_1(H_1^\dagger H_1)^2 + \frac{1}{2}\kappa_2(H_2^\dagger H_2)^2 \\ &+ \kappa_3(H_1^\dagger H_1)(H_2^\dagger H_2) + \kappa_4(H_1^\dagger H_2)(H_2^\dagger H_1) \\ &+ \left\{\frac{1}{2}\kappa_5(H_1^\dagger H_2)^2 + [\kappa_6(H_1^\dagger H_1) \right. \\ &+ \left. \kappa_7(H_2^\dagger H_2)]H_1^\dagger H_2 + \text{h.c.}\right\}. \end{aligned} \tag{7}$$

Allowing additional hard supersymmetry breaking quartic terms besides the usual gauge (D-)terms and loop contributions, $\kappa_{1\ldots7}$ can be written out explicitly as

$$\kappa_{1,2} = \frac{1}{2}(g'^2 + g^2) + \kappa_{soft\,1,2} + \kappa_{hard\,1,2}, \tag{8}$$

$$\kappa_3 = -\frac{1}{4}(g'^2 - g^2) + \kappa_{soft\,3} + \kappa_{hard\,3}, \tag{9}$$

$$\kappa_4 = -\frac{1}{2}g^2 + \kappa_{soft\,4} + \kappa_{hard\,4}, \tag{10}$$

$$\kappa_{5,6,7} = \kappa_{soft\,5,6,7} + \kappa_{hard\,5,6,7}, \tag{11}$$

where g' and g are the SM hypercharge and SU(2) gauge couplings, and $\kappa_{soft\,i}$ sums the loop effects due to soft supersymmetry breaking effects $\sim \ln m_0$. The effect of the HSB contributions $\kappa_{hard\,i}$ is estimated next.

While Ref. [8] explore the individual contribution of each of the $\kappa_{hard\,i=1,\cdots,7}$, here we will assume, for simplicity, that $\kappa_{hard\,i} = \kappa_{hard}$ are all equal and positive. The squared Higgs mass $m_{h^0}^2$ reads in this case

$$m_{h^0}^2 = M_Z^2\cos^2 2\beta + \delta m_{loop}^2 + (c_\beta + s_\beta)^4 v^2 \kappa_{hard}, \tag{12}$$

[5] In the decoupling limit it simply reduces to the SM with one "light" physical Higgs boson h^0, $m_{h^0}^2 = \kappa v^2$, $\kappa = c_\beta^4\kappa_1 + s_\beta^4\kappa_2 + 2s_\beta^2 c_\beta^2(\kappa_3 + \kappa_4 + \kappa_5) + 4c_\beta^3 s_\beta\kappa_6 + 4c_\beta s_\beta^3\kappa_7$, where $s_\beta \equiv \sin\beta$ and $c_\beta \equiv \cos\beta$, and $v = \langle h \rangle$.

Table 5
Frameworks for estimating κ_{hard}. (Saturation of the lower bound on M is assumed.)

	n	$\tilde{\kappa}_h$	M	$\delta\kappa_{hard}$
TLM	0	~ 1	$\gtrsim m_0$	$(m_0/M)^2 \sim 1$
NPGM	1/2	~ 1	$\gtrsim 4\pi m_0$	$(4\pi m_0/M)^2 \sim 1$
MGM	1	$\lesssim 1/16\pi^2$	$\gtrsim 16\pi^2 m_0$	$(4\pi m_0/M)^2 \sim 1/16\pi^2$

where $\delta m^2_{loop} \lesssim M_Z^2$ and $v = 174$ GeV is the SM Higgs VEV. (Note that no new particles or gauge interactions were introduced.)

Given the relation (12), one can evaluate the HSB contributions to the Higgs mass for an arbitrary M (and n). We define an effective scale $M_* \equiv (M/(4\pi)^{2n}\sqrt{\tilde{\kappa}_h})(\mathrm{TeV}/m_0)$. The HSB contributions decouple for $M_* \gg m_0$, and the results reduce to the MSSM limit with only SSB (e.g., supergravity mediation). However, for smaller values of M_* the Higgs mass is dramatically enhanced. For M=1 TeV and TLM or M=4π TeV and NPGM, both of which correspond to $M_* \simeq 1$ TeV, the Higgs mass could be as heavy as 475 GeV for $\tan\beta = 1.6$ and 290 GeV for $\tan\beta = 30$. This is to be compared with 104 GeV and 132 GeV [14], respectively, if HSB are either ignored or negligible. (The SM-like Higgs boson h^0 may be as heavy as 180 GeV in certain $U(1)'$ models [15] with only SSB.)

In the MGM case $\tilde{\kappa}_h \lesssim 1/16\pi^2$ so that $M_* \sim 4\pi$ TeV (unlike the NPGM where $M_* \sim 1$ TeV). HSB effects are now more moderate but can increase the Higgs mass by 40 (10) GeV for $\tan\beta = 1.6$ (30) (in comparison to the case with only SSB.) Although the increase in the Higgs mass in this case is not as large as in the TLM and NPGM cases, it is of the same order of magnitude as, or larger than, the two-loop corrections due to SSB [14], setting the uncertainty range on any such calculation.

In Fig. 1, m_{h^0} dependence on $\tan\beta$ for fixed values of M_* is shown. The $\tan\beta$ dependence is from the tree-level mass and from the HSB corrections, while the loop corrections to $m^2_{h^0}$ are fixed, for simplicity, at $9200\,\mathrm{GeV}^2$ [14]. The upper curve effectively corresponds to $\kappa_{hard} \simeq 1$. The HSB contribution dominates the Higgs mass and m_{h^0} decreases with increasing $\tan\beta$. As

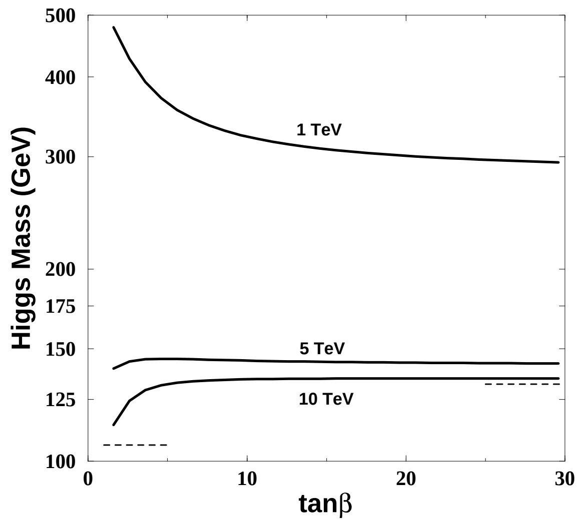

Figure 1. The light Higgs boson mass (note the logarithmic scale) is shown as a function of $\tan\beta$ for $M_* = 1, 5, 10$ TeV (assuming equal HSB couplings). The upper bound when considering only SSB ($M_* \to \infty$) is indicated for comparison (dashed lines) for $\tan\beta = 1.6$ (left) and 30 (right).

indicated above, m_{h^0} could be in the range of 300 − 500 GeV, dramatically departing from calculations which ignored HSB terms. The lower two curves illustrate the range[6] of the corrections in the MGM, where the tree-level and the HSB contributions compete. The $\cos 2\beta$ dependence of the tree-level term dominates the β-dependence of these two curves. Clearly, the Higgs mass could

[6] Given the many uncertainties, e.g., the messenger quantum numbers and multiplicity and $\sqrt{F}/M$ [10], we identify the MGM with a M_*-range which corresponds to a factor of two uncertainty in the hard coupling.

discriminate between the MGM and NPGM and help to better understand the origin of the supersymmetry breaking.

Following the Higgs boson discovery, it should be possible to extract information on the mediation scale M. In fact, some limits can already be extracted. Consider the upper bound on the Higgs mass derived from a fit to electroweak precision data: $m_h^0 < 215$ GeV at 95% confidence level [16]. (Such fits are valid in the decoupling limit discussed here.) A lower bound on the scale M in MGM could be obtained from

$$m_Z^2 \cos^2 2\beta + \delta m_{loop}^2 (c_\beta + s_\beta)^4 v^2 \left(\frac{4\pi m_0}{M}\right)^2 \leq (215 \text{ GeV})^2 \quad (13)$$

assuming equal κ_{hard}'s. For $\beta = 1.6$, it gives $M \geq 31$ TeV while for $\tan\beta = 30$ the lower bound is $M \geq 19$ TeV. Once m_{h^0} is measured, more stringent bounds on M could be set.

In conclusion, this section illustrates that the scale of the mediation of supersymmetry breaking explicitly appears in the prediction of the Higgs mass (and with a distinct β-dependence). In turn, it could lead in certain cases to a much heavier Higgs boson than usually anticipated in supersymmetric theories. It could also distinguish models, e.g., supergravity mediation from other low-energy mediation and weakly from strongly interacting messenger sectors. Given our ignorance of the (Kahler potential and) HSB terms, such effects can serve for setting the uncertainty on any Higgs mass calculations and can be used to qualitatively constrain the scale of mediation of supersymmetry breaking from the hidden to the SM sector.

4. SUMMARY

The most general treatment of the mediation of supersymmetry breaking requires the parameterization of the various SSB and HSB parameters in terms of the scale of supersymmetry breaking and of the scale of its mediation to the low-energy fields. This introduces in the infrared potential supersymmetry breaking parameters whose magnitude is not fixed by the weak scale, yet they do not destabilize the weak scale. While in some cases the parameters are very small, their effects need not be negligible. For example, in supergravity mediation they could lead to sizeable correction to the SSB parameters. In theories of low-energy supersymmetry breaking, such as gauge mediation, the effects could be even more dramatic, as illustrated in our discussion of the Higgs mass. Furthermore, the effects of, e.g., quartic couplings $\sim F^2/M^4$ may allow one to indirectly measure the scale M, as well as to determine the dynamics responsible for supersymmetry breaking (as in the case of strongly *vs.* weakly interacting messenger sectors).

Other interesting possibilities not explored here include chiral symmetry breaking and fermion mass generations [2]; theories of extended supersymmetry [2] and their accommodation of experimental constraints [17]; non-holomorphic trilinear interactions and their effect on the stability of the vacuum [6]; supergravity boundary conditions [7]; Higgsino mass generation [11]; theories with singlets [18]; and stabilization of flat potentials [3]. The relevance of various issues depends on the scale of supersymmetry breaking, which for the time being remains a free parameter. Treating it as such can illuminate some of the mysteries of the superworld as well as help in identifying various phenomena.

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